

CDM/baryon isocurvature perturbations in a sneutrino curvaton model

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Abstract

Matter isocurvature perturbations are strictly constrained from cosmic microwave background observations. We study a sneutrino curvaton model where both cold dark matter (CDM)/baryon isocurvature perturbations are generated. In our model, total matter isocurvature perturbations are reduced since the CDM/baryon isocurvature perturbations compensate for each other. We show that this model can not only avoid the stringent observational constraints but also suppress temperature anisotropies on large scales, which leads to improved agreement with observations.

1 Introduction

Inflation [1, 2, 3, 4, 5, 6, 7] in the very early universe solves horizon and flatness problems in the standard big bang cosmology. Furthermore, light scalar fields acquire fluctuations during inflation and can generate adiabatic and almost scale-invariant density perturbations which are in good agreement with cosmic microwave background (CMB) observations [8, 9, 10]. In the simplest case, a scalar field which causes inflation (called an inflaton) is responsible for the density perturbations, but it is also possible that another scalar field gives a significant contribution to the density perturbations. In particular, in curvaton models [11, 12, 13] some scalar field besides the inflaton (called a curvaton) obtains fluctuations during inflation and decays into radiation to produce adiabatic perturbations after reheating due to the inflaton decay.

However, in general, curvaton models produce not only adiabatic perturbations but also isocurvature ones. Isocurvature perturbations S_{ij} are defined as

$$S_{ij} = \frac{\delta\rho_i}{(1+w_i)\rho_i} - \frac{\delta\rho_j}{(1+w_j)\rho_j}, \quad (1)$$

where ρ_i , $\delta\rho_i$ and w_i are the energy density, its fluctuation and the coefficient of the equation of state of a component i . For example, if cold dark matter (CDM) (or baryon number) is generated and decouples from thermal bath before the curvaton decays, the isocurvature perturbations between CDM (baryon) and radiation are given by $S_{\text{CDM}(b)\gamma} = 0 - (3/4)(\delta\rho_\gamma/\rho_\gamma) \neq 0$, neglecting inflaton fluctuations. Moreover, the isocurvature perturbations produced in this way are anti-correlated with curvature perturbations ζ as $\zeta = -S_{\text{CDM}(b)\gamma}/3$. On the other hand, we have positively correlated isocurvature perturbations if CDM (baryon) is produced from the curvaton decay. In either case, since the isocurvature fluctuations are stringently constrained by CMB observations, they cause a serious cosmological difficulty in curvaton models.

An obvious solution to the isocurvature problem is to produce both CDM and baryon number thermally after the curvaton decays. Another interesting possibility is that CDM is produced from the inflaton and the baryon number is generated from the curvaton decay, or vice versa. Then, the total matter isocurvature perturbations S_m are written as

$$S_m = \frac{\Omega_{\text{CDM}}}{\Omega_m} S_{\text{CDM}\gamma} + \frac{\Omega_b}{\Omega_m} S_{b\gamma}, \quad (2)$$

where Ω_{CDM} , Ω_b and Ω_m are the density parameters of CDM, baryon and matter ($\Omega_m = \Omega_{\text{CDM}} + \Omega_b$), respectively. In this case, the contributions from CDM and baryon can cancel each other and hence we can avoid the stringent constraint from the CMB because the

temperature anisotropies are produced only through S_m (and ζ). When the cancellation occurs and S_m is significantly reduced in comparison with S_{CDM} and S_b , it is said that CDM and baryonic isocurvature perturbations are “compensated” with each other [14, 15, 16, 17, 18].

Furthermore, since the isocurvature perturbations can be anti-correlated with the adiabatic ones, temperature anisotropies on large scales can be reduced. In fact, the large scale temperature anisotropies by the Sachs-Wolfe effect is approximately given by [19]

$$\left\langle \left(\frac{\Delta T}{T} \right)^2 \right\rangle = \frac{1}{25} \left[\mathcal{P}_\zeta + 4\mathcal{P}_{S_m} + 4\mathcal{P}_{\zeta S_m} + \frac{5}{6}\mathcal{P}_T \right], \quad (3)$$

where \mathcal{P}_ζ and \mathcal{P}_{S_m} are the power spectra of curvature and isocurvature perturbations, and $\mathcal{P}_{\zeta S_m}$ is the cross power spectrum of them. Here we have included the contribution from tensor perturbations on large scales and \mathcal{P}_T is their power spectrum. From eq. (3), it is seen that $\Delta T/T$ decreases compared with a pure adiabatic case if $\mathcal{P}_{S_m} + \mathcal{P}_{\zeta S_m} < 0$. Moreover, when the tensor mode exits, the isocurvature perturbations can compensate for its effect on $\Delta T/T$. This effect may solve the tension between the Planck observation of the temperature fluctuations [10] and the BICEP2 detection of the B-mode polarization [20] as pointed out in [21]. Although it is premature to conclude the B-mode detection, taking into account uncertainty of the foreground dust emission [22, 23], need for reducing temperature fluctuations on large scales is also suggested from the analysis by [24], which shows that a negative tensor-to-scalar ratio $r \simeq -0.2$ gives a better fit to the Planck data. The compensated isocurvature perturbations can not only avoid the stringent CMB constraint but also have a possibility to improve agreement with the observational data.

In this paper we study a curvaton model where both CDM and baryonic isocurvature perturbations are produced. We identify the curvaton as a right-handed sneutrino in supersymmetric (SUSY) theories. In this model baryon asymmetry is generated via non-thermal leptogenesis by decay of sneutrinos into left-handed (s)leptons and higgs(inos). On the other hand, gravitinos produced during reheating after inflation can be dark matter if they are the lightest SUSY particles (LSPs). If not, LSPs produced by the gravitino decay account for dark matter. Then, baryon and CDM(=LSPs) have correlated and anti-correlated isocurvature perturbations, respectively, and they compensate for each other by taking appropriate model parameters. The possibility of such compensated isocurvature perturbations in a sneutrino curvaton model was pointed out in [25] but a detailed analysis has not been performed. We estimate the baryon number and the dark matter density as well as amplitudes of adiabatic and isocurvature perturbations in the sneutrino curvaton model. It is found that the compensated isocurvature perturbations are realized for the

reheating temperature of $O(10^{9-10})$ GeV and the LSP mass of $O(0.1-1)$ TeV.

This paper is organized as follows. In Sec. 2, we briefly explain how the CDM/baryon isocurvature perturbations could compensate for each other in the curvaton model. Then, we derive conditions for the CDM/baryon isocurvature perturbations to cancel each other. We also show conditions for the isocurvature perturbations to compensate for the contribution of the tensor mode to temperature anisotropies. In Sec. 3, we investigate a sneutrino curvaton model where the curvaton is identified as the right-handed sneutrino. We estimate perturbations generated by the curvaton, the baryon asymmetry, and the CDM abundance in our model. In Sec. 4, we express the above conditions for compensation with respect to model parameters of the curvaton scenario, such as the curvaton field value, the curvaton decay temperature and the mass of the LSP. Finally, in Sec. 5, we summarize our results.

2 Compensated isocurvature perturbations

In this section, we briefly review the curvaton scenario [11, 12, 13] and show how the isocurvature perturbations are generated. We explain how the compensated isocurvature perturbations could be realized in curvaton models.

In single-field inflation models, only the inflaton is responsible for the curvature perturbations. In curvaton scenarios, on the other hand, another light field called a curvaton is also the source of the adiabatic perturbations. During inflation, the curvaton is a subdominant component of the Universe. It acquires quantum fluctuations as the inflaton field does. After inflation ends, the inflaton begins its oscillation. Then, the curvaton field also starts to oscillate when the Hubble parameter becomes comparable to the curvaton mass. After that, the Universe is reheated via inflaton decay and a radiation dominated era is realized ^{#1}. During the radiation dominated period, the ratio of the curvaton density to the radiation energy density produced from the inflaton decay increases in proportional to a , where a is a scale factor. Thus, the curvaton becomes a non-negligible component of the Universe and the adiabatic perturbations evolve during this epoch. The evolution of the adiabatic perturbations stops when the curvaton decays and thereafter the curvature perturbations are conserved on super-horizon scales. The adiabatic perturbations in the curvaton scenario are given by

$$\zeta = \zeta_{\text{inf}} + \frac{f_{\text{dec}}}{3} S_{\sigma}, \quad (4)$$

^{#1} Here and hereafter, we assume that the curvaton starts to oscillate before the reheating since we identify the curvaton as the right-handed sneutrino whose mass is about $O(10^{12})$ GeV, which is heavy enough to begin oscillation before the reheating for typical reheating temperatures.

where ζ is the curvature perturbation on the uniform density slicing, ζ_{inf} is the curvature perturbation induced from the inflaton, and S_σ is the curvaton isocurvature perturbation which is given by $S_\sigma = 3(\zeta_\sigma - \zeta_{\text{inf}})$ with ζ_σ being the curvature perturbation on the uniform density slicing of the curvaton. The parameter f_{dec} is defined by the energy density of the radiation produced by the inflation ρ_r and that of the curvaton ρ_σ as

$$f_{\text{dec}} \equiv \frac{3\rho_\sigma}{(4\rho_r + 3\rho_\sigma)} \Big|_{\text{dec}}, \quad (5)$$

where the subscript “dec” denotes the value when the curvaton decays.

The curvaton scenario also generates matter isocurvature perturbations besides the adiabatic perturbations. As described in Sec. 1, the matter isocurvature perturbations consist of the CDM and the baryon isocurvature perturbations. If the CDM (baryon number) is produced from the decay products of the inflaton and decouples from thermal bath before the curvaton decays, the residual isocurvature perturbations are given by

$$\begin{aligned} S_{\text{CDM}(b)\gamma} &\equiv 3(\zeta_{\text{CDM}(b)} - \zeta) \\ &= -f_{\text{dec}} S_\sigma, \end{aligned} \quad (6)$$

with $\zeta_{\text{CDM}(b)} = \zeta_{\text{inf}}$. If the baryon number (CDM) is produced by out-of-equilibrium decay of the curvaton, the residual isocurvature perturbations are given by

$$S_{b(\text{CDM})\gamma} = (1 - f_{\text{dec}}) S_\sigma, \quad (7)$$

with $\zeta_{b(\text{CDM})} = \zeta_\sigma$. In this paper, we study the case where CDM is produced from decay products of the inflaton and the baryon number is produced by out-of-equilibrium decay of the curvaton. In this case, the adiabatic perturbations and the total matter isocurvature perturbations are given by

$$\begin{pmatrix} \zeta \\ S_m \end{pmatrix} = \begin{pmatrix} 1 & \frac{f_{\text{dec}}}{3} \\ 0 & \mathcal{T}_{S_m S_\sigma} \end{pmatrix} \begin{pmatrix} \zeta_{\text{inf}} \\ S_\sigma \end{pmatrix}, \quad (8)$$

where the transfer function $\mathcal{T}_{S_m S_\sigma}$ is given by

$$\begin{aligned} \mathcal{T}_{S_m S_\sigma} &= -\frac{\Omega_{\text{CDM}}}{\Omega_m} f_{\text{dec}} + \frac{\Omega_b}{\Omega_m} (1 - f_{\text{dec}}) \\ &= \frac{\Omega_b}{\Omega_m} - f_{\text{dec}}. \end{aligned} \quad (9)$$

Then, one finds that \mathcal{P}_ζ , \mathcal{P}_{S_m} and $\mathcal{P}_{\zeta S_m}$ are given by

$$\mathcal{P}_\zeta = \mathcal{P}_{\zeta_{\text{inf}}} + \frac{f_{\text{dec}}^2}{9} \mathcal{P}_{S_\sigma} \equiv (1 + R) \mathcal{P}_{\zeta_{\text{inf}}}, \quad (10)$$

$$\mathcal{P}_{S_m} = \mathcal{T}_{S_m S_\sigma}^2 \mathcal{P}_{S_\sigma}, \quad (11)$$

$$\mathcal{P}_{\zeta S_m} = \frac{f_{\text{dec}}}{3} \mathcal{T}_{S_m S_\sigma} \mathcal{P}_{S_\sigma}, \quad (12)$$

where R is defined as the ratio of the power spectra of the curvature perturbations from the inflaton to that from the curvaton. It should be noted that the cross power spectrum $\mathcal{P}_{\zeta S_m}$ is negative for $\mathcal{T}_{S_m S_\sigma} < 0$ and then anti-correlated isocurvature perturbations can be realized. It is found that the CDM/baryon isocurvature perturbations cancel each other and the total matter isocurvature perturbations vanish for $\mathcal{T}_{S_m S_\sigma} = 0$, in other words, $f_{\text{dec}} = \Omega_b/\Omega_m \simeq 0.16$ in eq. (9) for values of density parameter obtained from Planck+WP+highL+BAO [26].

Next let us focus on the temperature anisotropies. On large scales, the CMB anisotropies originate from the Sachs-Wolfe effect and are given by [19]

$$\left(\frac{\Delta T}{T} \right)_{\text{sw}} \simeq -\frac{1}{5} \zeta - \frac{2}{5} S_m + \frac{1}{2} h_{ij} n^i n^j, \quad (13)$$

where h_{ij} 's are the tensor perturbations and n is a unit vector. The correlation of the temperature anisotropies on large scales is approximately given by eq. (3),

$$\begin{aligned} \left\langle \left(\frac{\Delta T}{T} \right)^2 \right\rangle &= \frac{1}{25} \left[\mathcal{P}_\zeta + 4\mathcal{P}_{S_m} + 4\mathcal{P}_{\zeta S_m} + \frac{5}{6} \mathcal{P}_T \right] \\ &= \frac{1}{25} \mathcal{P}_\zeta \left[1 + 4B_m^2 + 4B_m \cos \theta_m + \frac{5}{6} r \right], \end{aligned} \quad (14)$$

where $r \equiv \mathcal{P}_T/\mathcal{P}_\zeta$. Here we have defined B_m and $\cos \theta_m$ as

$$B_m \equiv \sqrt{\frac{\mathcal{P}_{S_m}}{\mathcal{P}_\zeta}}, \quad (15)$$

$$\cos \theta_m \equiv \frac{\mathcal{P}_{\zeta S_m}}{\sqrt{\mathcal{P}_\zeta \mathcal{P}_{S_m}}}. \quad (16)$$

If the total matter isocurvature perturbations are anti-correlated with the adiabatic perturbations, the third term of eq. (14) is negative. In that case, it is possible that the isocurvature perturbations compensate for the tensor contributions to the temperature anisotropies on large scales [21]. In other words,

$$4\mathcal{P}_{S_m} + 4\mathcal{P}_{\zeta S_m} + \frac{5}{6} \mathcal{P}_T = 0, \quad (17)$$

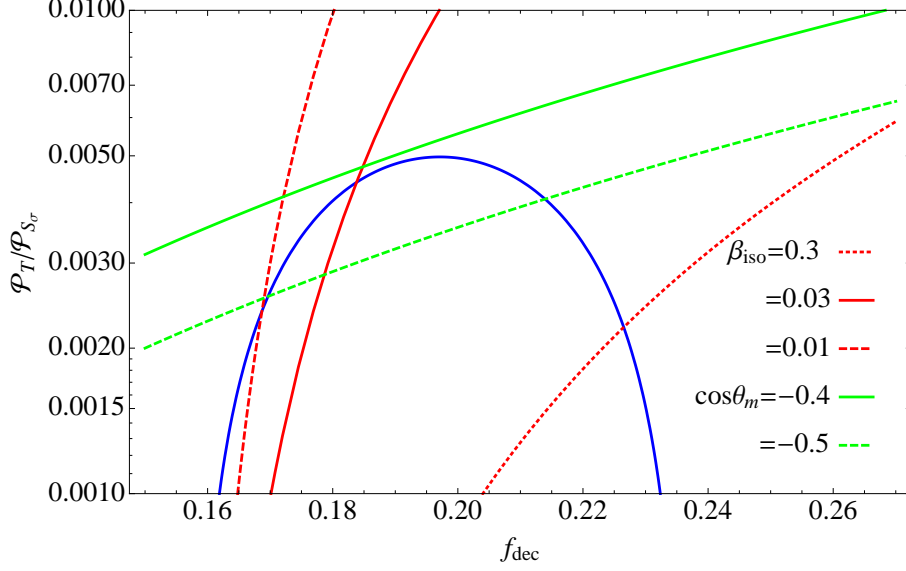


Figure 1: The solution for eq. (18) in $f_{\text{dec}}\text{-}\mathcal{P}_T/\mathcal{P}_{S_\sigma}$ plane. The red lines show contours of isocurvature fractions $\beta_{\text{iso}} \equiv \mathcal{P}_{S_m}/(\mathcal{P}_\zeta + \mathcal{P}_{S_m})$ and the green lines show those of $\cos\theta_m$ for $r = 0.2$.

could be satisfied. From eqs. (9), (11) and (12), we can rewrite it as follows,

$$4\mathcal{P}_{S_\sigma} \left(\frac{\Omega_b}{\Omega_m} - f_{\text{dec}} \right) \left(\frac{\Omega_b}{\Omega_m} - \frac{2}{3}f_{\text{dec}} \right) + \frac{5}{6}\mathcal{P}_T = 0. \quad (18)$$

The required isocurvature perturbations to satisfy eq. (18) are realized by taking appropriate parameters. In Fig. 1, we plot the solution of eq. (18) with the blue solid line in $f_{\text{dec}}\text{-}\mathcal{P}_T/\mathcal{P}_{S_\sigma}$ plane. We also plot contours of isocurvature fractions β_{iso} , defined as $\beta_{\text{iso}} \equiv \mathcal{P}_{S_m}/(\mathcal{P}_\zeta + \mathcal{P}_{S_m})$, and $\cos\theta_m$ defined in eq. (16), where we assume the tensor-to-scalar ratio $r = 0.2$. From observations of Planck and BICEP2 collaboration, the allowed parameter regions are $\beta_{\text{iso}} \lesssim 0.03$ and $\cos\theta_m \lesssim -0.4$ [27]. One can see that the solution exists around $f_{\text{dec}} \simeq 0.18$ and $\mathcal{P}_T/\mathcal{P}_{S_\sigma} \sim O(10^{-3})$. Hence, when we choose the parameters $\mathcal{P}_T/\mathcal{P}_{S_\sigma}$ and f_{dec} appropriately, we can compensate for the contribution from the tensor mode in the CMB temperature anisotropies on large scales, due to the anti-correlated isocurvature perturbations.

In this paper, we investigate following two cases:

- The CDM/baryon isocurvature perturbations cancel each other ($S_m = 0$).
- The isocurvature perturbations compensate for the tensor contribution to the temperature anisotropies.

As explained above, the former case is realized for $f_{\text{dec}} \simeq 0.16$ and the latter case is realized when the condition given by eq. (18) is satisfied.

3 Sneutrino curvaton scenario

In this section, we focus on a curvaton scenario based on supersymmetric theories. We identify the curvaton with a right-handed sneutrino in supersymmetric theories. We introduce three generations of right-handed neutrino chiral multiplets N_i with masses M_i ($i = 1, 2, 3$) into the minimal supersymmetric standard model in order to explain neutrino masses by the seesaw mechanism [28]. A superpotential of the right-handed neutrinos is given by

$$W = \frac{1}{2}M_i N_i N_i + \lambda_{ij} N_i L_j H_u, \quad (19)$$

where L_j and H_u are chiral multiplets of the lepton doublets and the up-type Higgs, respectively. We assume mass hierarchy as $M_1 \ll M_2 < M_3$ and that the lightest right-handed sneutrino \tilde{N}_1 plays the role of the curvaton #2.

We assume that the baryon asymmetry is generated via decay of the sneutrino curvaton. We also assume that gravitinos produced during reheating are responsible for the observed CDM abundance. In this way, baryon and CDM have correlated and anti-correlated isocurvature perturbations, respectively, and they compensate for each other by taking appropriate model parameters.

In the following, we express the power spectra and f_{dec} which are discussed in the previous section, in terms of model parameters of the sneutrino curvaton scenario. We also estimate the baryon asymmetry and the CDM abundance.

3.1 Power spectra and f_{dec}

We first express the power spectra in terms of model parameters. Hereafter, we assume that the curvaton has a quadratic potential, $V(\sigma) = \frac{1}{2}M_1^2\sigma^2$. The power spectrum of curvature perturbations generated from the inflaton fluctuations is given by

$$\mathcal{P}_{\zeta_{\text{inf}}}(k) = \frac{H_*^2(k)}{8\pi^2 M_{\text{pl}}^2 \epsilon_*(k)}, \quad (20)$$

where H is the Hubble parameter with the star denoting the epoch of horizon exit, $k = a_* H_*$, M_{pl} is the reduced Planck mass and ϵ_* is a slow-roll parameter defined as

#2 \tilde{N}_2 or \tilde{N}_3 can be the inflaton which causes chaotic inflation [29, 30].

$\epsilon_* \equiv \frac{M_{\text{pl}}^2}{2} \left(\frac{V_\phi}{V} \right)^2 \big|_{\phi_*}$ where the subscript ϕ denotes $\partial/\partial\phi$. The power spectrum of the curvaton isocurvature perturbations is given by [31]

$$\mathcal{P}_{S_\sigma} = \frac{4}{\sigma_*^2} \left(\frac{H_*}{2\pi} \right)^2, \quad (21)$$

where σ_* denotes the curvaton field value at the horizon exit. The power spectrum of the tensor perturbations is given by

$$\mathcal{P}_T = \frac{8}{M_{\text{pl}}^2} \left(\frac{H_*}{2\pi} \right)^2. \quad (22)$$

The ratio of the power spectra of the tensor to curvaton isocurvature perturbations is therefore given by

$$\frac{\mathcal{P}_T}{\mathcal{P}_{S_\sigma}} = 2 \left(\frac{\sigma_*}{M_{\text{pl}}} \right)^2. \quad (23)$$

One finds that $\mathcal{P}_T/\mathcal{P}_{S_\sigma}$ depends only on the curvaton field value σ_* .

Next we estimate the ratio of the curvaton to the radiation energy density at the curvaton decay time, $\rho_\sigma/\rho_r|_{\text{dec}}$, which appears in the parameter f_{dec} . Since the ratio of the energy densities does not change while both the inflaton and the curvaton are oscillating, it is given by

$$\frac{\rho_\sigma}{\rho_r} \bigg|_{\text{dec}} = \frac{\rho_\sigma}{\rho_\phi} \bigg|_{\text{osc}} \left(\frac{a_{\text{dec}}}{a_{\text{reh}}} \right), \quad (24)$$

where subscripts “osc”, “dec” and “reh” respectively denote the time of the onset of the curvaton oscillation, the curvaton decay and the reheating. We take $H_{\text{osc}} = M_1$, $H_{\text{dec}} = \Gamma_\sigma$ and $H_{\text{reh}} = \Gamma_\phi$, where Γ_σ and Γ_ϕ are decay rates of the curvaton and the inflaton. Then, we define the reheating (curvaton decay) temperature as $T_{\text{reh(dec)}} \equiv \left(\frac{90}{\pi^2 g_*} M_{\text{pl}}^2 \Gamma_{\phi(\sigma)}^2 \right)^{1/4}$. Assuming that the curvaton field stays at σ_* until its oscillation, $\rho_\sigma|_{\text{osc}} = \frac{1}{2} M_1^2 \sigma_*^2$. We also assume that the inflaton dominates the energy density of the Universe when the curvaton starts to oscillate. Then, $\rho_\sigma/\rho_\phi|_{\text{osc}}$ is written as

$$\begin{aligned} \frac{\rho_\sigma}{\rho_\phi} \bigg|_{\text{osc}} &= \frac{\frac{1}{2} M_1^2 \sigma_*^2}{3 M_{\text{pl}}^2 H_{\text{osc}}^2} \\ &= \frac{M_1^2 \sigma_*^2}{6 M_{\text{pl}}^2 M_1^2} = \frac{1}{6} \left(\frac{\sigma_*}{M_{\text{pl}}} \right)^2. \end{aligned} \quad (25)$$

In the second line, we have used $H_{\text{osc}} = M_1$. Since the temperature of the Universe is inversely proportional to the scale factor, eq. (24) is written as

$$\frac{\rho_\sigma}{\rho_r} \bigg|_{\text{dec}} = \frac{1}{6} \left(\frac{\sigma_*}{M_{\text{pl}}} \right)^2 \frac{T_{\text{reh}}}{T_{\text{dec}}}. \quad (26)$$

Since the parameter f_{dec} is determined by $\rho_\sigma/\rho_r|_{\text{dec}}$, it depends on the curvaton field value σ_* and the ratio of the temperature $T_{\text{reh}}/T_{\text{dec}}$.

Finally, we also estimate the parameter r defined in eq. (14). The tensor-to-scalar ratio r is given by

$$\begin{aligned} r &= \frac{\mathcal{P}_T}{(1+R)\mathcal{P}_{\zeta_{\text{inf}}}} \\ &= 16\epsilon_* \left[1 + \frac{8}{9}\epsilon_* f_{\text{dec}}^2 \left(\frac{M_{\text{pl}}}{\sigma_*} \right)^2 \right]^{-1}. \end{aligned} \quad (27)$$

From eqs. (23), (26) and (27), it is found that $\mathcal{P}_T/\mathcal{P}_{S_\sigma}$ depends only on σ_* , f_{dec} depends on σ_* and $T_{\text{reh}}/T_{\text{dec}}$, and r depends on σ_* , $T_{\text{reh}}/T_{\text{dec}}$ and ϵ_* .

3.2 Baryon asymmetry

\tilde{N}_1 non-thermally decays into (s)leptons and higgs(inos) as well as their anti-particles. Since CP symmetry is generally violated in the interaction of leptons, lepton asymmetry is generated by the decay of \tilde{N}_1 . The lepton asymmetry is partially converted into the baryon asymmetry through the sphaleron process [32].

The relation between the baryon asymmetry and the lepton asymmetry in supersymmetric theories is given by [33, 34]

$$\frac{n_B}{s} = -\frac{8}{23} \frac{n_L}{s}, \quad (28)$$

where s , n_B and n_L are respectively the entropy density, the baryon number density and the lepton number density. The lepton asymmetry generated by the decay of \tilde{N}_1 is given by

$$\frac{n_L}{s} = \epsilon_1 \frac{n_{\tilde{N}_1}}{s} \Big|_{\text{dec}}, \quad (29)$$

where the parameter ϵ_1 denotes CP asymmetry by the decay of \tilde{N}_1 , which is defined as

$$\epsilon_1 \equiv \frac{\Gamma(\tilde{N}_1 \rightarrow \tilde{L} + H_u) - \Gamma(\tilde{N}_1 \rightarrow \tilde{L}^* + H_u^*)}{\Gamma(\tilde{N}_1 \rightarrow \tilde{L} + H_u) + \Gamma(\tilde{N}_1 \rightarrow \tilde{L}^* + H_u^*)}, \quad (30)$$

where Γ 's are decay rates and tildes denote superpartners of the right-handed neutrinos and lepton doublets. Considering the interference of the one-loop diagrams with the tree level coupling in supersymmetric theories, the parameter ϵ_1 is given by [35]

$$\epsilon_1 = -\frac{1}{8\pi} \frac{1}{(\lambda\lambda^\dagger)_{11}} \sum_{i=2,3} \text{Im} \left[\{(\lambda\lambda^\dagger)_{1i}\}^2 \right] f \left(\frac{M_i^2}{M_1^2} \right), \quad (31)$$

where the function $f(x)$ is defined as

$$f(x) \equiv \sqrt{x} \ln \left(1 + \frac{1}{x} \right) + \frac{2\sqrt{x}}{x-1}. \quad (32)$$

Using the mass hierarchy $M_1 \ll M_2 < M_3$ and the seesaw relation, ϵ_1 is estimated as

$$\epsilon_1 \simeq \frac{3}{8\pi} \frac{M_1}{\langle H_u \rangle^2} m_{\nu 3} \delta_{\text{eff}}, \quad (33)$$

where $\langle H_u \rangle$ denotes the VEV of the up-type Higgs bosons, $m_{\nu 3}$ is the heaviest neutrino mass produced via the seesaw mechanism and δ_{eff} is a CP violating phase with $|\delta_{\text{eff}}| \leq 1$.

The number density of \tilde{N}_1 at its decay is given in terms of σ_* as

$$\begin{aligned} n_{\tilde{N}_1}|_{\text{dec}} &= \frac{1}{2} M_1 \sigma_*^2 \left(\frac{a_{\text{osc}}}{a_{\text{reh}}} \right)^3 \left(\frac{a_{\text{reh}}}{a_{\text{dec}}} \right)^3 \\ &= \frac{1}{2} M_1 \sigma_*^2 \left(\frac{\Gamma_\phi}{M_1} \right)^2 \left(\frac{T_{\text{dec}}}{T_{\text{reh}}} \right)^3. \end{aligned} \quad (34)$$

In the second line, we have used the following relation,

$$\left(\frac{a_{\text{osc}}}{a_{\text{reh}}} \right)^3 = \left(\frac{H_{\text{reh}}}{H_{\text{osc}}} \right)^2 = \left(\frac{\Gamma_\phi}{M_1} \right)^2. \quad (35)$$

By using the relations, $3M_{\text{pl}}^2 \Gamma_\phi^2 = \frac{\pi^2}{30} g_* T_{\text{reh}}^4$ and $s|_{\text{dec}} = \frac{2\pi^2}{45} g_{*s} T_{\text{dec}}^3$, we obtain

$$\frac{n_{\tilde{N}_1}}{s} \Big|_{\text{dec}} = \frac{1}{8} \left(\frac{\sigma_*}{M_{\text{pl}}} \right)^2 \frac{T_{\text{reh}}}{M_1}. \quad (36)$$

From eqs. (28) and (29), the generated baryon asymmetry by the decay of the sneutrino curvaton is given by

$$\begin{aligned} \frac{n_B}{s} &= -\frac{1}{23} \epsilon_1 \left(\frac{\sigma_*}{M_{\text{pl}}} \right)^2 \frac{T_{\text{reh}}}{M_1} \\ &\simeq -1.5 \times 10^{-11} \left(\frac{\sigma_*}{10^{17} \text{ GeV}} \right)^2 \left(\frac{T_{\text{reh}}}{10^9 \text{ GeV}} \right) \left\{ \left(\frac{174 \text{ GeV}}{\langle H_u \rangle} \right)^2 \left(\frac{m_{\nu 3}}{0.05 \text{ eV}} \right) \delta_{\text{eff}} \right\}. \end{aligned} \quad (37)$$

Hereafter, we take $\langle H_u \rangle \simeq 174 \text{ GeV}$ assuming $\langle H_u \rangle$ is larger than the VEV of the down-type Higgs, $\langle H_d \rangle$. As for the heaviest neutrino mass and the CP violating phase, we take $m_{\nu 3} = 0.05 \text{ eV}$ and $\delta_{\text{eff}} = -1$. We treat σ_* and T_{reh} as free parameters and choose them in order to obtain the present baryon asymmetry, $n_B/s \simeq 8.7 \times 10^{-11}$ [36].

3.3 CDM abundance

In supersymmetric theories, gravitinos are copiously produced during reheating. The gravitinos could be dark matter if they are the lightest SUSY particles (LSPs). If not, decay products of the gravitinos may account for dark matter. In the following, we assume that gravitinos produced during reheating is the dominant source of the dark matter abundance and that non-thermal production of gravitinos by decay of scalar condensations and thermal production of dark matter are negligible.

The gravitino abundance for the reheating temperature of $O(10^9)$ GeV is approximately given by [41]

$$Y_{3/2} \equiv \frac{n_{3/2}}{s} \simeq 1.4 \times 10^{-13} \left[1 + 0.6 \left(\frac{m_{1/2}}{m_{3/2}} \right)^2 \right] \left(\frac{T_{\text{reh}}}{10^9 \text{ GeV}} \right), \quad (38)$$

where $n_{3/2}$, $m_{3/2}$ and $m_{1/2}$ respectively denote the gravitino number density, the gravitino mass and the unified gaugino mass at the GUT scale.

From eq. (38), the dark matter density parameter is given by

$$\Omega_{\text{CDM}} h^2 \simeq 3.8 \times 10^{-2} \left[1 + 0.6 \left(\frac{m_{1/2}}{m_{3/2}} \right)^2 \right] \left(\frac{m_{\text{LSP}}}{1 \text{ TeV}} \right) \left(\frac{T_{\text{reh}}}{10^9 \text{ GeV}} \right). \quad (39)$$

In the case where the gravitino is not the LSP, assuming $m_{3/2}$ is much heavier than $m_{1/2}$, the second term of eq. (38) is negligible.

As mentioned in Sec. 3.2, we treat T_{reh} as a free parameter. The mass parameters, such as $m_{3/2}$, $m_{1/2}$ and m_{LSP} , depend on SUSY breaking models. Therefore, we also treat the mass parameters as free parameters. We choose T_{reh} and the mass parameters so as to obtain the correct dark matter density parameter, $\Omega_{\text{CDM}} h^2 \simeq 0.12$ [26].

4 Model parameters

In the previous section, we have investigated the generation of the CDM/baryon isocurvature perturbations and estimated the baryon asymmetry and the CDM abundance in the context of the sneutrino curvaton model. In this section, we investigate parameter regions in the model which realizes the compensated isocurvature perturbations and the case where the anti-correlated isocurvature perturbations compensate for the tensor contribution in the large scale CMB temperature anisotropies.

4.1 Compensation for isocurvature perturbations

We first consider the case where the compensation for the CDM/baryon isocurvature perturbations occurs, that is, $S_m = 0$, which requires $f_{\text{dec}} \simeq 0.16$. In this case, from eqs. (26) and (37) we find that the curvaton decay temperature T_{dec} is fixed to

$$T_{\text{dec}} \simeq 7 \times 10^6 \text{ GeV}. \quad (40)$$

This is a robust prediction in our model.

T_{reh} and σ_* depend on the mass parameters $m_{1/2}$, $m_{3/2}$ and m_{LSP} . In order to obtain the observed CDM density parameter, T_{reh} is given by

$$T_{\text{reh}} \simeq 3 \times 10^9 \text{ GeV} \left[1 + 0.6 \left(\frac{m_{1/2}}{m_{3/2}} \right)^2 \right]^{-1} \left(\frac{m_{\text{LSP}}}{1 \text{ TeV}} \right)^{-1}, \quad (41)$$

from eq. (39). From eq. (37) and the above value of T_{reh} , the condition for successful baryogenesis leads to

$$\sigma_* \simeq 1 \times 10^{17} \text{ GeV} \left[1 + 0.6 \left(\frac{m_{1/2}}{m_{3/2}} \right)^2 \right]^{1/2} \left(\frac{m_{\text{LSP}}}{1 \text{ TeV}} \right)^{1/2}. \quad (42)$$

The curvaton scenario is known as a model which may generate large local primordial non-Gaussianity. The non-linearity parameter $f_{\text{NL}}^{\text{loc}}$ has been commonly used as a parameter characterizing the local primordial non-Gaussianity. Planck collaboration gives a tight constraint on $f_{\text{NL}}^{\text{loc}}$ as $f_{\text{NL}}^{\text{loc}} < O(10)$ [37]. Basically, in the case where the contribution from the curvaton fluctuations is dominant in the power spectrum of the curvature perturbation, that is, $R \gg 1$, the non-linearity parameter is estimated as $f_{\text{NL}}^{\text{loc}} \sim 1/f_{\text{dec}}$. On the other hand, in the case with $R \ll 1$, $f_{\text{NL}}^{\text{loc}}$ is suppressed as $f_{\text{NL}}^{\text{loc}} \sim R^2 \times 1/f_{\text{dec}}$. In our model, R is given by $R = \frac{8}{9} \epsilon_* f_{\text{dec}}^2 \left(\frac{M_{\text{pl}}}{\sigma_*} \right)^2$. Therefore, the parameter region we consider corresponds to the case with $R \sim O(0.1)$, and hence our model predicts $f_{\text{NL}}^{\text{loc}} = O(0.1)$, which is consistent with the Planck result.

It is found that the curvaton field value σ_* is slightly smaller than M_{pl} and the reheating temperature T_{reh} is of $O(10^{9-10})$ GeV, where we assume $m_{\text{LSP}} \sim m_{1/2} \sim O(1)$ TeV and $m_{3/2} \gtrsim O(1)$ TeV. The reheating temperature of $O(10^{9-10})$ GeV is a typical one in chaotic inflation models [38, 39, 40]. If the gravitino is the LSP and much lighter than the weak scale, σ_* is far above M_{pl} , which is incompatible with the curvaton scenario. We stress that the curvaton decay temperature T_{dec} is uniquely predicted as $T_{\text{dec}} \simeq 7 \times 10^6$ GeV in this model.

4.2 Compensation for tensor modes

Next, we consider the case where the anti-correlated isocurvature perturbations compensate for the tensor modes. In Sec. 2, we have derived the condition, eq. (18), on $\mathcal{P}_T/\mathcal{P}_{S_\sigma}$ and f_{dec} to realize the compensation for tensor modes. In Sec. 3.1, we have shown the dependence of $\mathcal{P}_T/\mathcal{P}_{S_\sigma}$ and f_{dec} on model parameters, σ_* and $T_{\text{reh}}/T_{\text{dec}}$. In Fig. 2, we show the conditions to realize the compensation for tensor modes with the blue solid line in σ_* - $T_{\text{reh}}/T_{\text{dec}}$ plane, corresponding to Fig. 1. Note that they have no dependence on the tensor-to-scalar ratio r . The black lines show the condition to obtain r for each slow-roll parameter ϵ_* , estimated by eq. (27). The red lines are contours of the isocurvature fractions $\beta_{\text{iso}} \equiv \mathcal{P}_{S_m}/(\mathcal{P}_\zeta + \mathcal{P}_{S_m})$ and the green lines are those of $\cos\theta_m$.

As mentioned in Sec. 2, constraints on β_{iso} and $\cos\theta_m$ are given by $\beta_{\text{iso}} \lesssim 0.03$ and $\cos\theta_m \lesssim -0.4$ [27]. Furthermore, since the observed curvature perturbations are nearly scale invariant and the spectral index is $n_s \simeq 0.96$ [10], the slow-roll parameter should be small as $\epsilon_* \lesssim O(0.01)$. Therefore, we find that

$$\sigma_*/M_{\text{pl}} \sim O(0.01) \quad \text{and} \quad T_{\text{reh}}/T_{\text{dec}} \sim O(10^3), \quad (43)$$

are required in order to realize the compensation for tensor modes.

In Fig. 3, we plot the relation between the curvaton field value σ_* and the reheating temperature T_{reh} to obtain the present baryon asymmetry. From eq. (43) and Fig. 3, we find that $T_{\text{reh}} = O(10^9)$ GeV is required.

Once the allowed values of the reheating temperature T_{reh} and the curvaton field value σ_* are confined, the mass parameters $m_{3/2}$, $m_{1/2}$ and m_{LSP} are determined from eq. (39) to obtain the observed dark matter abundance. We find that

$$m_{3/2} \gtrsim O(0.1\text{-}1) \text{ TeV}, \quad m_{1/2} \sim m_{\text{LSP}} \sim O(0.1\text{-}1) \text{ TeV}, \quad (44)$$

are required if the gravitino is as heavy as or heavier than minimal supersymmetric standard model particles. It is remarkable that required soft masses are consistent with TeV scale SUSY. If the gravitino is the LSP and much lighter than the weak scale, $m_{1/2}$ is required to be smaller, which is inconsistent with various experiments.

As we have mentioned in Sec. 4.1, the curvaton model might generate large primordial non-Gaussianity, $f_{\text{NL}}^{\text{loc}}$. The parameter region we consider, however, corresponds to the case with $R \sim O(0.1)$ and hence our model is consistent with the Planck result.

As we have shown, our model realizes the isocurvature perturbations which compensate for the tensor modes and reduce the temperature fluctuations on large scales, which can solve the possible tension between the Planck observation of the temperature fluctuations [10] and the BICEP2 detection of the B-mode polarization [20], as pointed out

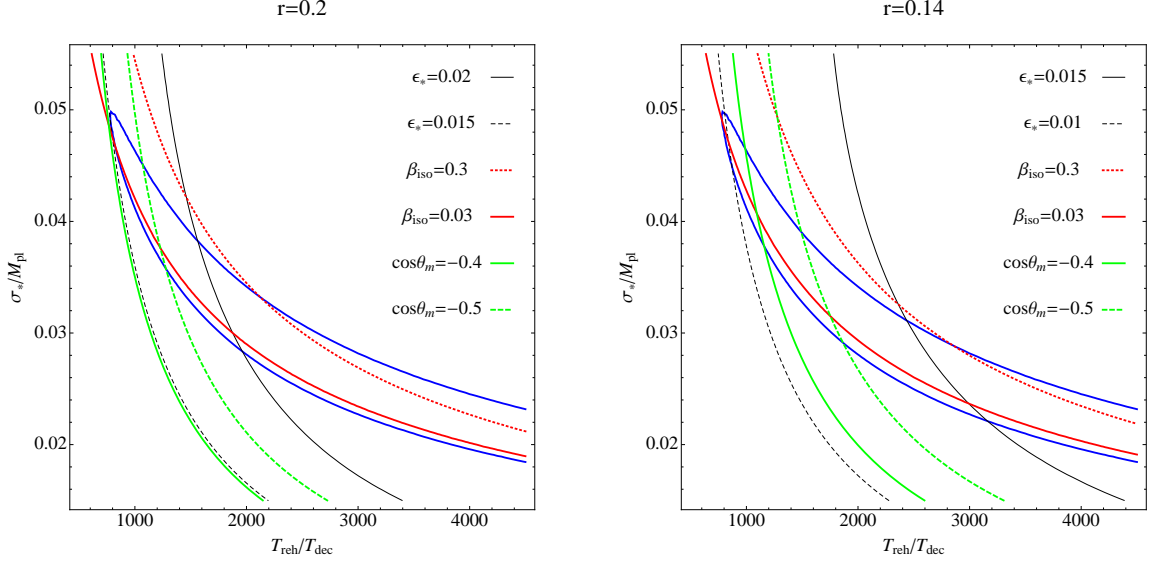


Figure 2: The solutions for eq. (18) in σ_* - $T_{\text{reh}}/T_{\text{dec}}$ plane are denoted by the blue lines. The black lines show the condition to obtain the tensor-to-scalar ratio r for each slow-roll parameter ϵ_* . The red lines are contours of the isocurvature fractions β_{iso} and the green lines are those of $\cos \theta_m$. The allowed parameter region, $\beta_{\text{iso}} < 0.03$ and $\cos \theta_m < -0.4$, is below the red solid line and above the green solid line. We show the cases for $r = 0.2$ (left panel) and $r = 0.14$ (right panel).

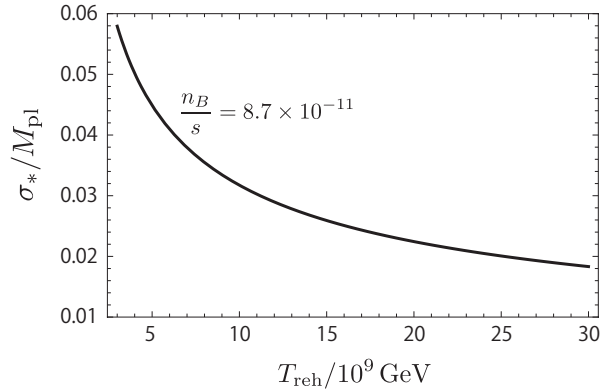


Figure 3: The present baryon asymmetry in σ_* - T_{reh} plane.

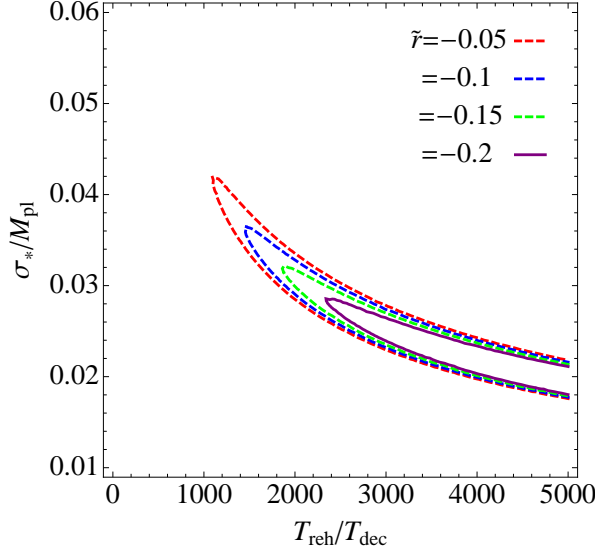


Figure 4: The solutions for eq. (45) in σ_* - T_{reh} plane. We fix the slow-roll parameter $\epsilon_* = 0.01$.

in [21]. We also note that a negative tensor-to-scalar ratio $r \simeq -0.2$ may give a better fit to the Planck data [24]. From eq. (14), we can realize such a situation when the following relation is satisfied for $\tilde{r} = -0.2$,

$$4 \frac{\mathcal{P}_{S_m}}{\mathcal{P}_\zeta} + 4 \frac{\mathcal{P}_{\zeta S_m}}{\mathcal{P}_\zeta} + \frac{5}{6} r = \frac{5}{6} \tilde{r}, \quad (45)$$

where \tilde{r} is an effective tensor-to-scalar ratio. The left-hand side depends on σ_* , $T_{\text{reh}}/T_{\text{dec}}$ and ϵ_* . We plot the condition to realize the negative tensor-to-scalar ratio for each value of \tilde{r} in Fig 4, where we assume $\epsilon_* = 0.01$.

5 Summary and Discussions

In this paper, we have investigated the sneutrino curvaton model in which the CDM/baryon isocurvature perturbations are generated. We have considered cases where the isocurvature perturbations compensate each other or compensate for the tensor contributions to the temperature anisotropies, which could improve agreement between observations. In our curvaton scenario, the non-thermal leptogenesis from the sneutrino curvaton is responsible for the baryon asymmetry and gravitinos produced during the reheating or their decay products account for dark matter. We have found that the compensation requires

the curvaton field value during inflation of $O(10^{17})$ GeV, the reheating temperature T_{reh} of $O(10^{9-10})$ GeV and the curvaton decay temperature T_{dec} of $O(10^{6-7})$ GeV.

Finally, let us comment on mass spectra of SUSY particles, based on the constraint from the BBN [41, 42, 43]. When the gravitino is the LSP, the gravitino mass is of $O(0.1-1)$ TeV. Decay of the next-to LSP (NLSP) might destroy the success of the BBN. If the NLSP is the left-handed sneutrino, however, such a problem could be avoided since the left-handed sneutrino dominantly decays into particles whose interactions are weak [41].

On the other hand, when the LSP is not the gravitino, the gravitino may be far heavier than the LSP. If the gravitino mass is larger than $O(10)$ TeV, the constraint from the BBN is avoided. To obtain the observed dark matter density, $m_{\text{LSP}} \simeq O(0.1-1)$ TeV $\ll m_{3/2}$ is required. Such hierarchy is naturally explained if gaugino mass is generated only by the anomaly mediation [44, 45], as is the case with high scale SUSY breaking models [46, 47, 48, 49].

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